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Critique of a Pion Exchange Model for Interquark Forces

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I describe four serious defects of a widely discussed pion exchange model for interquark forces: it doesn't solve the "spin-orbit problem" as advertised, it fails to describe the internal structure of baryon resonances, it leads to disastrous conclusions when extended to mesons, and it is not reasonably connected to the physics of heavy-light systems.

I. INTRODUCTION

The idea that the low-energy degrees of freedom of QCD are quarks, gluons, and Goldstone bosons is an old and interesting one. In one form or another, it has been used in a wide variety of models for the last two decades [1,2]. In this paper I offer a critique of a recent and widely discussed variant of such models due to Glozman and Riska [1] in which it is proposed that baryon spectroscopy be described by discarding the standard one-gluon-exchange (OGE) forces of De Rújula, Georgi, and Glashow [3] (which were applied to baryons most extensively by Isgur and Karl [4]) in favor of the exchange between quarks of the octet of pseudoscalar mesons (OPE). I will avoid the distraction of criticizing either practical details of the Glozman-Riska model [1] or its theoretical foundations. I will instead accept the original model at face value and describe what I see as its four most serious defects. Some of these defects have been pointed out less formally in the past, and partly in response the original OPE model [1] has been elaborated [5–8]. Although these elaborations have not overcome the defects I describe here [9], I will comment upon them as appropriate in what follows.

II. A CATALOGUE OF CRITICISMS

A. The Spin-Orbit Problem is Not Solved

One of the central motivations for the Glozman-Riska model was to solve the "baryon spin-orbit problem". The Isgur-Karl model [4] discards spin-orbit forces in view of the data which demand that such forces be small, so that many of the successes of that model are due to the OGE-induced hyperfine interactions (of both the spin-spin and tensor types). The authors of Ref. [1] note that OPE produces hyperfine interactions *without* spin-orbit interactions, and argue that this supports their hypothesis that OPE is the true origin of the residual interquark forces (*i.e.,* of interactions beyond those which produce confinement).

This argument has a fundamental flaw. The zeroth-order confining potential, whose eigenstates are the basis for first order perturbation theory in both the Isgur-Karl and Glozman-Riska models, will produce very strong spin-orbit forces through Thomas precession, a purely kinematic effect. From the observed spectrum of states, it is impossible to escape the conclusion that this source of spin-orbit forces alone would produce *inverted* spin-orbit multiplets with splittings of *hundreds of MeV*. Thus the true nature of the "spin-orbit problem" seems to have been misunderstood: it is to arrange a sufficiently precise *cancellation* between dynamically generated spin-orbit forces and the inevitable Thomas-precession-induced spin-orbit forces.

These issues are discussed in the original Isgur-Karl papers [4], but especially in view of some recent developments in the subject, I will review the main points here. For reasons that will soon become apparent, it is best to start the discussion in the meson sector. The mesons also have a "spin-orbit problem" as can be seen by examining the first band of positive parity excited mesons: the four P-wave mesons of every flavor are nearly degenerate. Most of the observed small non-degeneracies are due to hyperfine interactions, but the spin-orbit matrix elements can be extracted. For example, by taking the isovector meson combination $\frac{5}{12}m_{a_2} - \frac{1}{4}m_{a_1} - \frac{1}{6}m_{a_0}$ one can isolate their spin-orbit matrix element of -3 ± 20 MeV. As in the baryons, this matrix element is much smaller than would be obtained from OGE. However, as already explained, this is not the point. In fact, a substantial "normal" spin-orbit matrix element is *needed* to cancel the strong "inverted" spin-orbit matrix element from Thomas precession in the confining potential [10]. For example, a recent fit [11] to the data on heavy-light mesons, including as a limiting case the light-light isovector mesons, gives an OGE spin-orbit matrix element of $+240$ MeV and a Thomas precession spin-orbit matrix element from the confining potential of -200 MeV: both are very large but they are nearly perfectly cancelling.

The physics behind this cancellation has received support recently from analyses of heavy quarkonia, where both analytic techniques [12] and numerical studies using lattice QCD [13] have shown that the confining forces are spin-independent *apart* from the inevitable spin-

orbit pseudoforce due to Thomas precession. Moreover, as has been known for more than ten years, the data on charmonia require a negative spin-orbit matrix element from Thomas precession in the confining potential to cancel part of the strength of the positive OGE matrix element. If the charm quark were sufficiently massive, its low-lying spectrum would be rigorously dominated by one gluon exchange. Indeed, one observes that the Υ system is closer to this ideal, as expected. Conversely, as one moves from $c\bar{c}$ to lighter quarks, the $\ell = 1$ wave functions move farther out into the confining potential and the relative strength of the Thomas precession term grows. It is thus very natural to expect a strong cancellation in light quark systems, though the observed nearly perfect cancellation must be viewed as accidental.

With these points in mind, let us now turn to baryons. As shown in the original Isgur-Karl paper on the P-wave baryons [4], a very similar cancellation can occur at the two-body level in baryons. However, unlike mesons, baryons can also experience three-body spin orbit forces [14] (*e.g.,* potentials proportional to $(\vec{S}_1 - \vec{S}_2) \cdot (\vec{r}_1 - \vec{r}_2) \times \vec{p}_3$ where $\vec{S}_i$, $\vec{r}_i$, $\vec{p}_i$ are the spin, position, and momentum of quark i). The matrix elements of these three body spin-orbit forces are all calculated in Ref. [4], but no apparent cancellation amongst them is found. *I.e.,* the spin-orbit problem might more properly be called the "baryon three-body spin orbit problem". In view of the facts that one could understand the smallness of spin-orbit forces in mesons and that the data clearly called for small spin-orbit forces in baryons, the Isgur-Karl model anticipated a solution to the baryon three-body spin-orbit problem and *as a first approximation* discarded all spin-orbit forces. It was assumed that, as in mesons, a more precise and broadly applicable description would have to treat residual spin-orbit interactions [15].

It should now be clear that replacing OGE by OPE is *not* a step forward, but rather a step backward, in dealing with the observed smallness of spin-orbit forces in baryons. By eliminating the OGE spin-orbit forces the Glozman-Riska model has not solved the baryon three-body spin-orbit problem but it has fully exposed (*i.e.,* left completely uncancelled) the strong Thomas precession forces from confinement. Thus, this model has escalated the

baryon spin-orbit problem into a "baryon two- and three-body spin-orbit problem", not solved it. The recent extensions of the Glozman-Riska model to include vector meson exchange [5,8] have attempted to address this problem. It is found that ρ exchange can produce a very strong spin-orbit interaction which might cancel with Thomas precession. Needless to say, at the best such an arrangement offers no improvement over the original situation [16]! Since, for reasons to be described below, it would require independent solutions of the meson and baryon spin-orbit puzzles, I consider this aspect of even the elaborated Glozman-Riska model to be a decisive step backward.

B. Baryon Internal Wave Functions are Wrong

In a complex system like the baryon resonances, predicting the spectrum of states is not a very stringent test of a model. The prototypical example (and the first case in N^* spectroscopy where this issue arises) is the two $N^*\frac{1}{2}^-$ states found in the 1500-1700 MeV range. In any reasonable valence quark model, two $N^*\frac{1}{2}^-$ states will be predicted in this mass range: the excitation of a unit of orbital angular momentum will create the negative parity and cost about 500 MeV in excitation energy relative to the $N - \Delta$ center-of-mass position at 1100 MeV (c.f. the $a_2 - \rho$ splitting), and totally antisymmetric states with overall angular momentum $\frac{1}{2}$ can be formed by coupling either quark spin $\frac{3}{2}$ or quark spin $\frac{1}{2}$ with $\ell = 1$. In the general case such a model will therefore give

$$|N^*\frac{1}{2}^- (upper)\rangle = cos\theta_{\frac{1}{2}^-}|^4P_N\rangle + sin\theta_{\frac{1}{2}^-}|^2P_N\rangle \tag{1}$$

$$|N^*\frac{1}{2}^- (lower)\rangle = cos\theta_{\frac{1}{2}^-}|^2P_N\rangle - sin\theta_{\frac{1}{2}^-}|^4P_N\rangle \tag{2}$$

in an obvious notation. Since the masses of these resonances are only known (and currently interpretable) to roughly 50 MeV, it is not extremely difficult to arrange for a model to give a satisfactory description of the $N^*\frac{1}{2}^-$ spectrum. However, among models which perfectly describe the spectrum there is still a continuous infinity of predictions for the internal composition of these two states since all values of $\theta_{\frac{1}{2}^-}$ from 0 to π correspond to distinct

states.

It had been appreciated for some time [17] that the peculiar decay properties of the $N^* \frac{1}{2}^-$ states, and in particular the dominance of the $N\eta$ decay of the lower state despite its phase space suppression relative to $N\pi$, required that $\theta_{\frac{1}{2}^-} \simeq -35°$. One of the early successes of the Isgur-Karl model was that it makes the *parameter free* prediction $\theta_{\frac{1}{2}^-} = -arctan(\frac{\sqrt{5}-1}{2}) \simeq -32°$! The quark model also predicts a pair of $N^* \frac{3}{2}^-$ states which have an analogous mixing angle $\theta_{\frac{3}{2}^-}$. The empirically determined value of that angle was $\theta_{\frac{3}{2}^-} \simeq +10°$, while the Isgur-Karl model predicts $\theta_{\frac{3}{2}^-} = arctan(\frac{\sqrt{10}}{14+\sqrt{206}}) \simeq +6°$. In contrast, though the OPE model produces a very acceptable negative parity N^* spectrum, it predicts $\theta_{\frac{1}{2}^-} = \pm 13°$ and $\theta_{\frac{3}{2}^-} = \pm 8°$. Even though Ref. [1] only quotes the probabilities of $|^4P_N\rangle$ admixtures so that the critical signs of these mixing angles are not available, these results are sufficient for one to see that *the internal structure of the predicted states is wrong.* In concrete terms, such a $\theta_{\frac{1}{2}^-}$, even if it has the right sign, will have almost no impact on explaining the anomalously large $N\eta$ branching ratio of the $N^*(1535)\frac{1}{2}^-$ and the anomalously small $N\eta$ branching ratio of the $N^*(1650)\frac{1}{2}^-$.

More extreme cases of the importance of using the internal structure of states and not just spectroscopy as tests of dynamics are found in the positive parity band of excited baryons in the 1700-2000 MeV range. For example, the valence quark model predicts *five* $N^* \frac{3}{2}^+$ states in this range, but only one is known. Given that the masses of the N^*'s are rarely known to better than 50 MeV, it would be an unlucky modeller who couldn't identify one of their five predicted states with the observed state and claim spectroscopic success! What is far less trivial, as in the $N^* \frac{1}{2}^-$ sector, is to ask whether the one "predicted" $N^* \frac{3}{2}^+$ state has production and decay amplitudes consistent with the observed state, and, equally important, to understand why the other four $N^* \frac{3}{2}^+$ states "did not bark in the night". This is part of the well-known missing resonance problem and, as shown in Ref. [18], the OGE mechanism of the Isgur-Karl model provides a remarkably complete explanation across the entire baryon spectrum for which states should have been seen *and* where they are seen [19]. There is no evidence that the OPE model has this critical property.

Our discussion of the definitive role of internal structure would be incomplete without an example which touches back on the issue of spin-orbit forces. That the $\Lambda(1405)\frac{1}{2}^-$ and $\Lambda(1520)\frac{3}{2}^-$ are not degenerate seems to be a sign that the approximation of neglecting spin-orbit forces is imperfect. Indeed, the discrepancy between the Isgur-Karl model prediction of 1490 MeV and the observed mass of 1405 MeV for the lightest $\Lambda\frac{1}{2}^-$ state is one of the model's worst spectroscopic failures. This has led to speculation that the $\Lambda(1405)\frac{1}{2}^-$ is a $\bar{K}N$ bound state. However, there is little doubt that, while its mass is off by 85 MeV, the predicted state is to be identified with the $\Lambda(1405)\frac{1}{2}^-$. An analysis [17] of the production and decay amplitudes of the three expected $\Lambda\frac{1}{2}^-$ baryons gives a best fit with

$$|\Lambda(1405)\frac{1}{2}^-\rangle = +0.80|^2\Lambda_1\rangle + 0.60|^2\Lambda_8\rangle - 0.04|^4\Lambda_8\rangle \tag{3}$$

$$|\Lambda(1670)\frac{1}{2}^-\rangle = -0.44|^2\Lambda_1\rangle + 0.63|^2\Lambda_8\rangle + 0.64|^4\Lambda_8\rangle \tag{4}$$

$$|\Lambda(1775)\frac{1}{2}^-\rangle = +0.41|^2\Lambda_1\rangle - 0.49|^2\Lambda_8\rangle + 0.77|^4\Lambda_8\rangle \tag{5}$$

where $^2\Lambda_1$ is the quark spin $\frac{1}{2}$ SU(3) singlet Λ and $^2\Lambda_8$ and $^4\Lambda_8$ are the quark spin $\frac{1}{2}$ and $\frac{3}{2}$ SU(3) octet Λ's, respectively. The Isgur-Karl model gives (again with no parameters)

$$|\Lambda(1490)\frac{1}{2}^-\rangle = +0.90|^2\Lambda_1\rangle + 0.43|^2\Lambda_8\rangle + 0.06|^4\Lambda_8\rangle \tag{6}$$

$$|\Lambda(1650)\frac{1}{2}^-\rangle = -0.39|^2\Lambda_1\rangle + 0.75|^2\Lambda_8\rangle + 0.53|^4\Lambda_8\rangle \tag{7}$$

$$|\Lambda(1800)\frac{1}{2}^-\rangle = +0.18|^2\Lambda_1\rangle - 0.50|^2\Lambda_8\rangle + 0.85|^4\Lambda_8\rangle \tag{8}$$

which is imperfect, but quite acceptable given the uncertainties in the data and in its interpretation.

Similarly, the decay analyses also indicate that the $\Lambda(1520)\frac{3}{2}^-$ is to be identified with the lightest $\Lambda\frac{3}{2}^-$ of the Isgur-Karl model. Experiment therefore tells us that, despite the spectroscopic discrepancies, the $\Lambda(1405)\frac{1}{2}^-$ and $\Lambda(1520)\frac{3}{2}^-$ are indeed spin-orbit partners which will evolve (as m_s increases to m_c and then to the heavy quark limit $m_Q = \infty$) into the degenerate partners of a heavy quark symmetry spin multiplet [20]. This fact will play an important role in Section D below.

It has been suggested [7] that the extension of the Glozman-Riska model to include other meson exchanges will correct the failure of the OPE model to describe the internal structure of the baryon resonances. This may be, but it remains to be demonstrated. It has also been claimed that recent phenomenological analyses of the negative parity baryons [21,22] are incompatible with the Isgur-Karl model, but support the Glozman-Riska model. Since the Isgur-Karl model fits the data reasonably well, this can hardly be true! In fact, for the reasons described in this Section, it is clear that the phenomenological matrix elements deduced by these analyses must be inconsistent with OPE. What both analyses do show is that a *generic* flavor-exchange interaction can produce a better fit to the data (with its *experimental* errors) than a generic flavor-independent model. Given the physics that is currently being ignored in such models (i.e., the unknown *theoretical* errors), the significance of this observation is unclear.

C. Mesons Are a Disaster

There are two ways in which the OPE model is a disaster for mesons: it doesn't produce spin-dependent interactions where they are needed and so requires that we invoke *independent* mechanisms for creating splittings in mesons and baryons, and it predicts the existence of effects in mesons which are ruled out experimentally.

We know from quenched lattice QCD that at least the bulk of both meson and baryon hyperfine interactions occur in the quenched approximation (*i.e.,* in the absence of closed $q\bar{q}$ loops) [23]. Figures 1 show a Z-graph-induced meson exchange between quarks that arises in the quenched approximation and could therefore in principle be the origin of the OPE-induced hyperfine interactions posited in the Glozman-Riska model. The first problem I wish to highlight is that this mechanism can only operate between two quarks and *not* between a quark and an antiquark, so if baryon spin-dependent interactions are dominated by OPE, meson and baryon spin-dependent interactions must have totally different physical origins. This is not only unaesthetic: as we shall see, it is also very difficult to arrange.
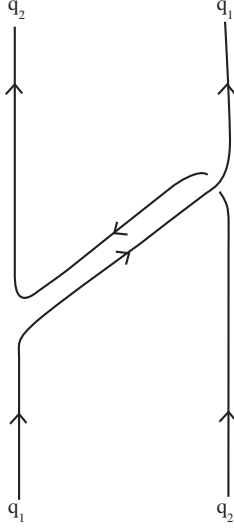
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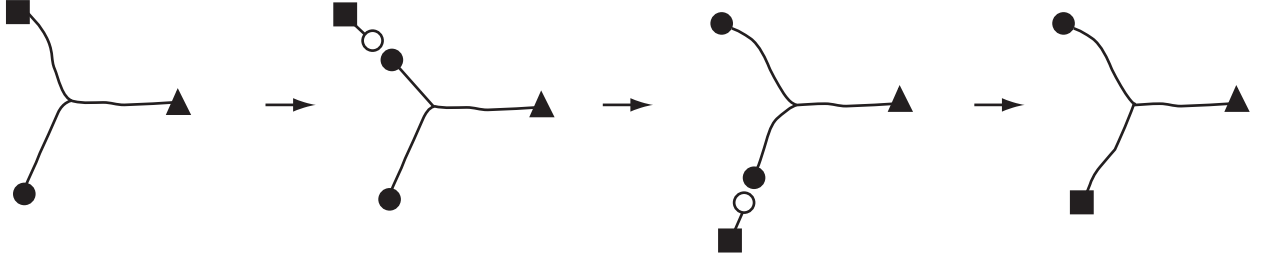Fig. 1(a): Z-graph-induced meson exchange between two quarks.

Fig. 1(b): A cartoon of the space-time development of the Z-graph-induced meson exchange in a baryon in the flux tube model. For diagrammatic clarity three different flavors of quarks are shown. Note that if the created meson rejoins the flux tube from which it originated, the produced $q\bar{q}$ pair can be of any flavor; however, such a process would be a closed $q\bar{q}$ loop and therefore not part of the quenched approximation. Also possible, but not shown, are OZI-violating graphs with the creation or annihilation of a disconnected $q\bar{q}$ meson; these are irrelevant to octet meson exchange in the SU(3) limit and enter in broken SU(3) only through the $\eta - \eta'$ mixing angle.

Figure 2 shows what we know about the evolution of quarkonium spectroscopy as a function of the quark masses. In heavy quarkonia ($b\bar{b}$ and $c\bar{c}$) we know that hyperfine interactions are generated by one-gluon-exchange perturbations of wave functions which are solutions of the Coulomb-plus-linear potential problem. I find it difficult to look at this diagram and not see a smooth evolution of the wavefunction (characterized by the slow evolution of the orbital excitation energy) convoluted with the predicted $1/m_Q^2$ strength of the OGE hyperfine interaction.
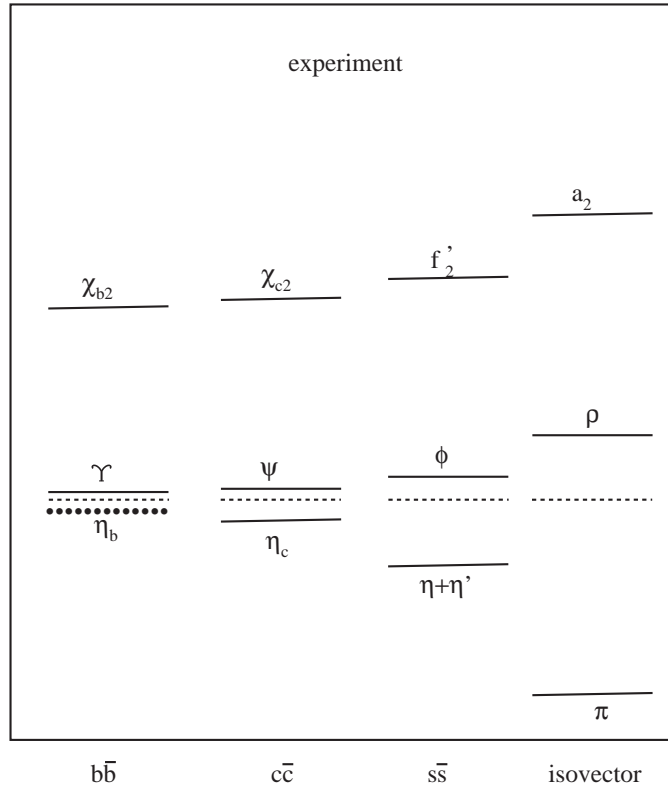

Fig. 2: The experimental spectra of $b\bar{b}$, $c\bar{c}$, $s\bar{s}$, and isovector light quarkonia, with the center of gravity of the S-wave mesons aligned. The 2^{++} states have been used to represent the P-wave mesons. The pseudoscalar $s\bar{s}$ state ("$\eta + \eta'$") has been located by unmixing a 2×2 matrix assumed to consist of primordial $s\bar{s}$ and $\frac{1}{\sqrt{2}}(u\bar{u} + d\bar{d})$ states. The η_b is not yet discovered, but the theoretical prediction is shown as a dotted spectral line. The spectra are shown to scale, which may conveniently be calibrated with the $\chi_{c2} - \psi$ splitting of 459 MeV.
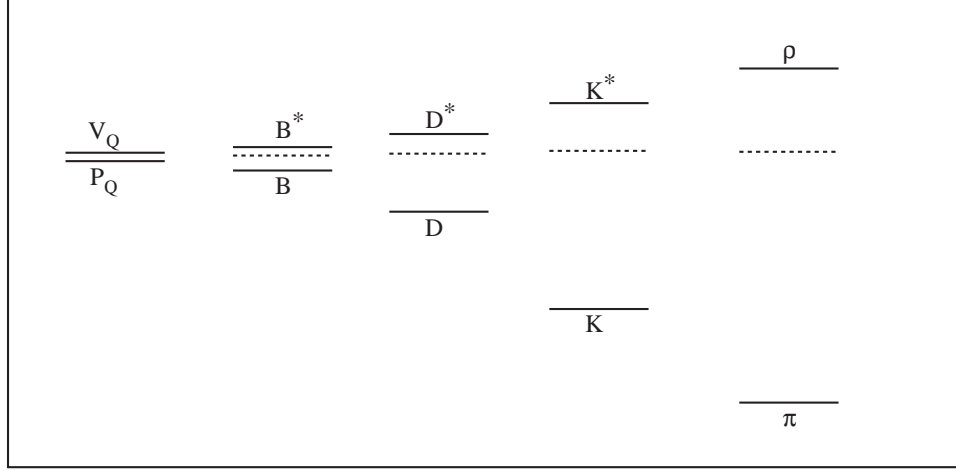
(a)
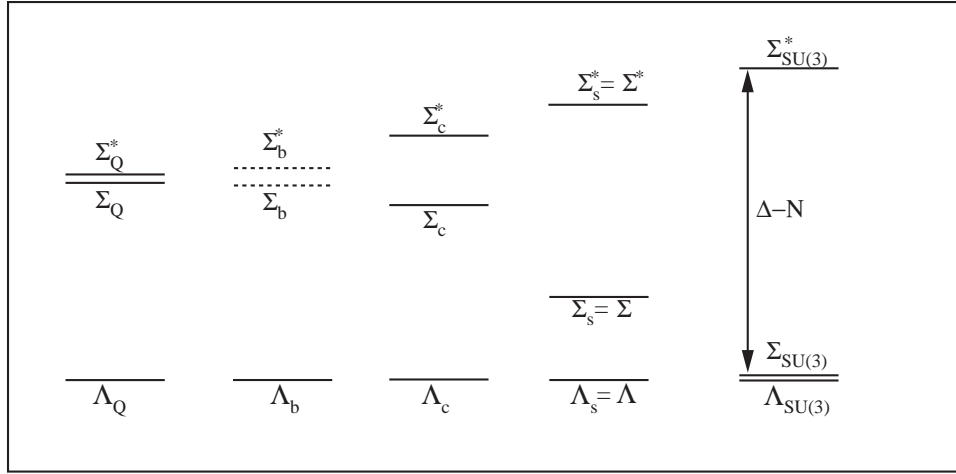
(b)

Fig. 3: Ground state meson (a) and baryon (b) hyperfine splittings in heavy-light systems as a function of the mass m_Q of the heavy quark. The spectra on the far left are the $m_Q \to \infty$ limits of heavy quark symmetry. The $\Sigma_Q^* - \Lambda_Q$ splitting and the positions of Σ_b^* and Σ_b are estimates from the quark model; all other masses are from experiment. The spectra are shown to scale; the meson scale may conveniently be calibrated with the $D^* - D$ splitting of 141 MeV and the baryon scale with the $\Sigma_c - \Lambda_c$ splitting of 169 MeV.

This same conclusion can be reached by approaching the light quarkonia from another angle. Figure 3(a) shows the evolution of heavy-light meson hyperfine interactions from the heavy quark limit to the same isovector quarkonia. In this case we know that in the heavy

quark limit [20] the hyperfine interaction is given by the matrix element of the operator $\vec{\sigma}_Q \cdot \vec{B}/2m_Q$. In contrast to heavy quarkonium, however, we do not know that the chromomagnetic field $\vec{B}$ at the position of Q is being generated by one gluon exchange from a light valence antiquark. Nevertheless, by considering unequal mass heavy quarkonia $Q\bar{q}$ with m_q beginning at m_Q and decreasing to the light quark mass m_d, one finds that the OGE hyperfine interaction extrapolates very neatly from the end of the region where it may be rigorously applied ($m_q \simeq 1$ GeV) down to light quark masses. The conclusion that heavy-light meson hyperfine interactions are controlled by OGE is also supported by the striking $1/m_Q$ behaviour of the ground state splittings in Fig. 3(a) as m_Q is decreased from m_b to m_c to m_s to m_d: it certainly appears that for all quark masses the quark Q interacts with $\vec{B}$ through its chromomagnetic moment $\vec{\sigma}_Q/2m_Q$, as would be characteristic of the OGE mechanism [24].

Since the OPE mechanism cannot contribute in mesons, the OGE mechanism is thus the natural candidate for generating meson hyperfine interactions. The objective (as opposed to aesthetic) problem that arises for the OPE hypothesis is that it is then nearly impossible to avoid the conclusion that OGE is also dominant in baryon hyperfine interactions: the OGE $q\bar{q}$ and qq hyperfine interactions are related by a simple factor of 1/2, and given the similarities of meson and baryon structure (for example, their charge radii, orbital excitation energies, and magnetic moments are all similar), it is inevitable that the matrix elements of OGE in baryons and mesons are similar. Valence quark model calculations support this qualitative argument, finding that mesons and baryons can be described by a universal confining potential with one-gluon exchange at short distances [25].

There is another very serious problem with the OPE mechanism which surfaces in mesons. I have explained that there are no Z-graph-induced meson exchanges in mesons. However, Fig. 4 shows how the same meson exchanges which are assumed to exist in baryons will drive mixings in isoscalar channels *by annihilation graphs*. More mechanically, the OPE mechanism posits the existence of vertices by which quarks couple to pseudoscalar mesons; antiquarks necessarily couple with the same strength to the charge conjugate mesons. By

considering the flavor structure of the allowed vertices, it is easy to show that the resulting pseudoscalar meson exchange between the quark and antiquark in a meson must have the character shown in Fig. 4, *i.e.*, it can only operate in isoscalar channels.
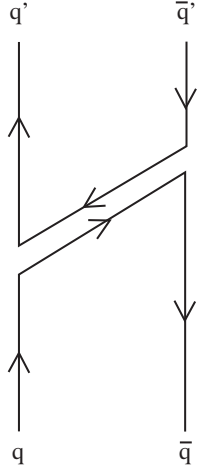

Fig. 4: OZI-violating mixing in isoscalar mesons via the exchange of a $q\bar{q}'$ meson.

I have argued above that the structure of mesons and baryons is so similar that it is impossible to avoid their having similar OGE matrix elements. The same is true for OPE matrix elements: it is impossible to maintain that OPE is strong enough to produce the $\Delta - N$ splitting in baryons without predicting a matrix element of comparable strength associated with Fig. 4 in mesons. Such matrix elements will violate the OZI rule [26]. Consider the mixing between the pure ω-like state $\frac{1}{\sqrt{2}}(u\bar{u} + d\bar{d})$ and the pure ϕ-like state $s\bar{s}$. This mixing will be driven by kaon exchange and from the preceding very general arguments we must expect that the amplitude A_{OZI} for this OZI-violating process will have a strength of the same order as the 200 MeV $\Sigma^* - \Sigma$ splitting (which is also driven purely by kaon exchange). Such an amplitude would be an order of magnitude larger than that observed: A_{OZI} for the vector mesons is very tiny - - - of the order of 10 MeV - - - corresponding to the known near purity of the ϕ as an $s\bar{s}$ state. The only escape from this disaster is to argue for some mechanism external to the model which could cancel the large kaon exchange contribution to A_{OZI}. Given that A_{OZI} is of the order of 10 MeV in not only the 1^{--} mesons,

but also in all of the other known meson nonets (except the pseudoscalars), this escape route seems implausible.

The mesons thus produce some disastrous conclusions for the Glozman-Riska model. The first is the very unaesthetic conclusion that two totally distinct mechanisms are in operation producing meson and baryon spin-dependent interactions: OGE in mesons and OPE in baryons. The second is the virtual impossibility of having strong OGE matrix elements in mesons without also producing strong OGE matrix elements in baryons, in conflict with the basic hypothesis of that model. The third is that the OPE mechanism produces unacceptably large OZI violation in meson nonets.

D. The Connection to Heavy Quark Baryons is Lost

As shown in Fig. 3(b), the baryon analog of Fig. 3(a), experiment provides further strong evidence in support of the dominance of OGE and *not* OPE in the baryons themselves! It is clear from this Figure that in the heavy quark limit the OPE mechanism is not dominant: exchange of the heavy pseudoscalar meson P_Q would produce a hyperfine interaction that scales with heavy quark mass like $1/m_Q^2$, while for heavy-light baryons the splittings are behaving like $1/m_Q$ as in the heavy-light mesons. This is as demanded by heavy quark theory where these splittings are once again rigorously controlled by the matrix elements of $\vec{\sigma}_Q \cdot \vec{B}/2m_Q$.

It is difficult to look at this diagram and not see a smooth evolution of this $1/m_Q$ behaviour from m_c to m_s to m_d, where by SU(3) symmetry $\Sigma^*_{SU(3)} - \Lambda_{SU(3)} = \Delta - N$, the splitting under discussion here. Indeed, using standard constituent quark masses ($m_d = m_u \equiv m = 0.33$ GeV, $m_s = 0.55$ GeV, $m_c = 1.82$ GeV, $m_b = 5.20$ GeV), the OGE mechanism with its natural $1/m_Q$ behaviour *quantitatively* describes these spectra. In this picture, Σ^*, Σ, and Λ are the analogs of the heavy quark states Σ^*_Q, Σ_Q, and Λ_Q. I speak of analogs here because the heavy quark expansion cannot be justified for such light values of m_Q. Nevertheless, one expects and observes in both mesons and baryons the remnants

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or analogs of heavy quark spectroscopy in light quark systems. For example, in Fig. 3(a) the $D^* - D$ heavy quark spin multiplet is naturally identified with the $K^* - K$ multiplet, *i.e.,* the basic degrees of freedom seen in the spectrum are the same, and from the observed $D^* - D$ splitting and the $1/m_Q$ heavy quark scaling law one expects a $K^* - K$ splitting of 460 MeV, quite close to the actual splitting of 400 MeV. One might try to escape this conclusion by arguing that between m_c and m_s the OGE-driven $1/m_Q$ mechanism turns off in baryons and the $1/m_Q^2$ OPE mechanism turns on. From the baryon spectra alone, one cannot rule out this baroque possibility. However, in the heavy-light mesons of Fig. 3(a) there is no alternative to the OGE mechanism, and if the $Q\bar{q}$ interaction continues to grow like $1/m_Q$ as m_Q gets lighter, then (given the similarity of meson and baryon structure) so must the Qq interaction. I see no escape from the conclusion that OGE is dominant in *all* ground state hyperfine interactions.

The excited charmed baryon sector has recently provided further strong evidence for the dominance of the OGE mechanism in baryons. Recall the conclusion of Section B above that the $\Lambda(1405)\frac{1}{2}^-$ and $\Lambda(1520)\frac{3}{2}^-$ are spin-orbit partners. Heavy quark symmetry [20] demands that in the heavy-light isospin zero Λ_Q sector, the $\Lambda_Q\frac{1}{2}^-$ and $\Lambda_Q\frac{3}{2}^-$ be degenerate as $m_Q \to \infty$ and that their splitting open up like $1/m_Q$ as m_Q decreases. The $\Lambda_c(2594)\frac{1}{2}^-$ and $\Lambda_c(2627)\frac{3}{2}^-$ [27] appear to be such a nearly degenerate pair of states in the charmed baryon sector. The center-of-gravity $\frac{2}{3}m_{\Lambda_Q\frac{3}{2}^-} + \frac{1}{3}m_{\Lambda_Q\frac{1}{2}^-}$ of these two states is 330 MeV above the $\Lambda_c(2285)$. This is to be compared with the center-of-gravity of the $\Lambda(1520)$ and $\Lambda(1405)$ which lies 365 MeV above the $\Lambda(1115)$, in accord with the expectation from the quark model that the orbital excitation energy of the negative parity excitations of Λ_Q will be a slowly increasing function of $1/m_Q$. (The quark model makes a similar prediction for the P-wave heavy-light mesons which is confirmed by the data.) This alone suggests that the strange quark analogs of the heavy quark spin multiplet $(\Lambda_Q\frac{3}{2}^-, \Lambda_Q\frac{1}{2}^-)$ which $(\Lambda_c(2627)\frac{3}{2}^-, \Lambda_c(2594)\frac{1}{2}^-)$ exemplifies should exist just around the mass of the $\Lambda(1520)\frac{3}{2}^-$ and $\Lambda(1405)\frac{1}{2}^-$. Moreover, using the predicted $1/m_Q$ behaviour of the $(\Lambda_Q\frac{3}{2}^-, \Lambda_Q\frac{1}{2}^-)$ multiplet splitting would lead to a predicted splitting of 110 MeV in the Λ sector compared to the observed splitting of 115 MeV. It is

thus very difficult to avoid identifying $(\Lambda(1520)\frac{3}{2}^-, \Lambda(1405)\frac{1}{2}^-)$ as the strange quark analogs of a heavy quark spin multiplet, and to avoid concluding that the $1/m_Q$ evolution of the OGE mechanism is responsible for its splitting.

Heavy quark symmetry thus poses another serious problem for the Glozman-Riska model. The OPE mechanism not only fails to explain meson hyperfine interactions, but it also cannot even explain the spectra of all baryons: it violates the requirement that splittings in heavy-light baryons open up like $1/m_Q$. Heavy-light mesons show that the $1/m_Q$ behaviour of the OGE mechanism persists all the way down to light quark masses, and the observed behaviour of baryons is completely consistent with the same extension of the heavy quark symmetry scaling laws. This behaviour leaves very little room for the OPE mechanism, and certainly makes it implausible that it is dominant.

III. CONCLUSIONS

I have focused in this paper on the predictions of the Glozman-Riska OPE model [1]. I believe the catalogue of problems I have described are sufficient to rule the model out.

Though in its extreme version the model is unsustainable, some elements of the physics of the Glozman-Riska model surely play a role in baryons. While the process depicted in Fig. 1(a) may be taken into account in a "dual approximation" by relativistic valence quark propagators, quark-antiquark correlations in the intermediate state will produce departures from this approximation. We may expect that such departures will be most pronounced in situations where the meson spectrum most strongly breaks the closure limit required for duality [28–30], and the pion certainly has the potential to do this. Of course, given that these exchanges occur at short distances and that the structure of the Goldstone bosons is very similar to that of other mesons [31], there is no obvious rationale for truncating the tower of meson exchanges associated with Fig. 1 with these states alone. (I hasten to add that it is *not* the distance to threshold which directly determines the importance of a given meson: its dominant contribution comes from the peak of its spectral function, and

in realistic models this feature is controlled more by the internal structure of a meson than by its mass [28–30].)

The $1/N_c$ expansion offers additional insights into this issue and into the structure of the arguments presented in this paper. Since OGE and OPE are of the same order in $1/N_c$ in baryons, there is unlikely to be any general principle which could be used to decide which is dominant. This situation can be contrasted with the case of heavy quarks for which the Z-graphs of Fig. 1(a) are suppressed like Λ_{QCD}/m_Q relative to gluon exchange graphs. The dominance of OGE over OPE in baryons must therefore have a dynamical origin. In contrast to baryons, OPE-like effects in mesons are suppressed by a power of $1/N_c$ relative to gluon exchange. One can thus look to mesons for a relatively unobscured picture of the strength and character of OGE effects, as I did in many of the arguments of this paper.

It remains to speculate on why $q\bar{q}$ pair creation effects are not more important in baryons. The OPE (or more generally meson-exchange) potentials between quarks would only be the simplest manifestation of such effects. More generally, meson emission and reabsorption by the baryons would lead to a complex interaction of the discrete baryon spectrum with each of the baryon-meson continua. Studies of "unquenching the quark model" have provided a plausible explanation for why such effects do not demolish the spectroscopy of the quark potential models [28,29] and the success of the OZI rule [30]. These studies indicate that the resiliency of valence quark model spectroscopy to $q\bar{q}$ pair creation occurs *not* because such processes are intrinsically weak, but because most of their effects can be absorbed into renormalized valence quark model parameters. The prime example of this absorption of $q\bar{q}$ effects is the string tension. Though the string tension is a strong function of the number of light flavors, when it is carefully renormalized to its observed value it produces the observed spectrum. A particularly treacherous aspect of this situation is that if one examines the effect of any particular continuum channel on the spectrum, it may appear to be very large and to have a complex dependence on the state of the quarks; only the sum over all channels leads in first approximation to a simple renormalization of the string tension. After this summation, only small residual effects associated with nearby thresholds

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remain [28].

The renormalization of the string tension is associated with $q\bar{q}$ pair creation in which the pair forms a closed loop. However, a created $q\bar{q}$ pair of the appropriate flavor could also join with a valence quark or antiquark to make a Z-graph. As already mentioned, the sum over all such processes at the hadronic level is apparently dual to the leading Z-graph component of the relativistic valence quark propagator. In baryons, meson exchange between quarks is contained in the Z-graph process of Fig. 1(a), and therefore the bulk of the effect of such hadronic processes should be absorbed into renormalized valence quark model parameters, in this case those describing the relativistic valence quark propagator. I speculate that duality is once again sufficiently accurate that only a small residue of this second type of $q\bar{q}$ effect remains after summing over channels, suppressing meson exchange relative to its naive strength in a $1/N_c$ expansion.

It would be interesting to apply the methods of Refs. [28–30] to the Z-graphs to see if the effects of duality violation are indeed small. However, whatever the reason, the empirically-based arguments of this paper make it clear that gluon exchange and not meson exchange provides the dominant residual forces between constituent quarks in both mesons and baryons.

ACKNOWLEDGEMENTS

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relativistically from a point of meson creation to a point of annihilation. That there is nothing particularly special about the Goldstone bosons is supported by the fact that their physical size as measured by their form factors is very similar to those of any other meson (*e.g.,* the $\bar{B}$, the D, the $\bar{K}$, and the π all have very similar charge radii). Their excitation spectra also seem to be totally normal. These observations immediately raise some issues which must be addressed in certain OPE models. In a formulation where one takes as basic degrees of freedom quarks, gluons, and Goldstone bosons, there will be a double-counting problem in the meson sector since it will have both a "fundamental" Goldstone boson and a quark-antiquark bound state Goldstone boson. The second problem that must be faced is that it is not legitimate to treat the quark-Goldstone boson vertex as pointlike: this vertex will be quite soft because the Goldstone bosons are quite large. Thus while the approximation of using a point-like quark-Goldstone boson coupling is appropriate at large distances where chiral perturbation theory applies, a cutoff will have to be applied for the small interquark distances inside a proton, and results will typically be strongly dependent on this cutoff.